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A Realistic World from Intersecting D6-Branes

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We briefly describe a three-family intersecting D6-brane model in Type IIA theory on the $\mathbf{T^6}/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold with a realistic phenomenology. In this model, the gauge symmetry can be broken down to the Standard Model (SM) gauge symmetry close to the string scale, and the gauge coupling unification can be achieved. We calculate the supersymmetry breaking soft terms, and the corresponding low energy supersymmetric particle spectrum, which may be tested at the Large Hadron Collider (LHC). The observed dark matter density may also be generated. Finally, we can explain the SM quark masses and CKM mixings, and the tau lepton mass. The neutrino masses and mixings may be generated via the seesaw mechanism as well.

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Introduction – During the last few years, intersecting D-brane models on Type II orientifolds [1], where the chiral fermions arise from the intersections of D-branes in the internal space [2] and the T-dual description in terms of magnetized D-branes [3] have been particularly interesting [4]. On Type IIA orientifolds with intersecting D6-branes, a large number of non-supersymmetric three-family Standard-like models and Grand Unified Theories (GUTs) were constructed in early stages [5]. However, there generically existed uncancelled Neveu-Schwarz-Neveu-Schwarz tadpoles in these models as well as the gauge hierarchy problem. To solve these problems, semi-realistic supersymmetric Standard-like and GUT models have been constructed in Type IIA theory on the $\mathbf{T^6}/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold [6, 7] and other backgrounds [8]. To stabilize the moduli via supergravity fluxes, flux models on Type II orientifolds have also been constructed [9, 10]. There are two main constraints on supersymmetric D-brane model building: (1) Ramond-Ramond (RR) tadpole cancellation conditions and (2) the requirement for four-dimensional $N = 1$ supersymmetric D-brane configurations.

However, there are two serious problems in almost all supersymmetric D-brane models: the absence of gauge coupling unification at the string scale, and the rank one problem in the Standard Model (SM) fermion Yukawa matrices. Thus, a comprehensive phenomenological study of a concrete model from the string scale to the electroweak scale has yet to be made. Although these problems can be solved in the flux models of Ref. [10] where the RR tadpole cancellation conditions are relaxed, those models are in the AdS vacua and the resulting flux induced superpotential for moduli is too complicated. Interestingly, we find that there is one and only one intersecting D6-brane model on the Type IIA $\mathbf{T^6}/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold where the above problems can be solved [7, 10]. Therefore, it is desirable to study the phenomenological consequences of this model in great detail.

Model Building – We consider Type IIA string theory compactified on a $\mathbf{T^6}/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold [6]. The $\mathbf{T^6}$ is a six-torus factorized as $\mathbf{T^6} = \mathbf{T^2} \times \mathbf{T^2} \times \mathbf{T^2}$ whose complex coordinates are z_i, $i = 1,\ 2,\ 3$ for the i^{th} two torus, respectively. The θ and ω generators for the orbifold group $\mathbb{Z}_2 \times \mathbb{Z}_2$, act on the complex coordinates of $\mathbf{T^6}$ as

$$\begin{aligned} \theta: &\ (z_1, z_2, z_3) \to (-z_1, -z_2, z_3)\ , \\ \omega: &\ (z_1, z_2, z_3) \to (z_1, -z_2, -z_3)\ . \end{aligned} \quad (1)$$

The orientifold projection is implemented by gauging the symmetry ΩR, where Ω is world-sheet parity, and R is given by

$$R: (z_1, z_2, z_3) \to (\overline{z}_1, \overline{z}_2, \overline{z}_3)\ . \quad (2)$$

Thus, there are four kinds of orientifold 6-planes (O6-planes) for the actions ΩR, $\Omega R\theta$, $\Omega R\omega$, and $\Omega R\theta\omega$, respectively. There are two kinds of complex structures consistent with orientifold projection for a two torus: rectangular and tilted [6]. If we denote the homology classes of the three cycles wrapped by the D6-brane stacks as $n^i_P[a_i] + m^i_P[b_i]$ and $n^i_P[a'_i] + m^i_P[b_i]$ with $[a'_i] = [a_i] + \frac{1}{2}[b_i]$ for the rectangular and tilted tori respectively, we can label a generic one cycle by (n^i_P, l^i_P) in either case, where in terms of the wrapping numbers $l^i_P \equiv m^i_P$ for a rectangular two torus and $l^i_P \equiv 2\tilde{m}^i_P = 2m^i_P + n^i_P$ for a tilted two torus. Moreover, for a stack of N D6-branes that does not lie on one of the O6-planes, we obtain a $U(N/2)$ gauge symmetry with three adjoint chiral superfields due to the orbifold projections, while for a stack of N D6-branes which lies on an O6-plane, we obtain a $USp(N)$ gauge symmetry with three anti-symmetric chiral superfields. Bifundamental chiral superfields arise from the intersections of two different stacks P and Q of D6-branes or from one stack P and its ΩR image P' [6].

We present the D6-brane configurations and intersection numbers of the model in Table I, and the resulting spectrum in Table II [7, 10]. We put the a', b, and c stacks of D6-branes on the top of each other on the third two torus, and as a resut there are additional vector-like particles from $N=2$ subsectors.

TABLE I: D6-brane configurations and intersection numbers.

	$U(4)_C \times U(2)_L \times U(2)_R \times USp(2)^4$												
	N	$(n^1,l^1)\times(n^2,l^2)\times(n^3,l^3)$	n_S	n_A	b	b'	c	c'	1	2	3	4
a	8	$(0,-1)\times(1,1)\times(1,1)$	0	0	3	0	-3	0	1	-1	0	0
b	4	$(3,1)\times(1,0)\times(1,-1)$	2	-2	-	-	0	0	0	1	0	-3
c	4	$(3,-1)\times(0,1)\times(1,-1)$	-2	2	-	-	-	-	-1	0	3	0
1	2	$(1,0)\times(1,0)\times(2,0)$	$\chi_1=3,\ \chi_2=1,\ \chi_3=2$									
2	2	$(1,0)\times(0,-1)\times(0,2)$	$\beta^g_1=-3,\ \beta^g_2=-3$									
3	2	$(0,-1)\times(1,0)\times(0,2)$	$\beta^g_3=-3,\ \beta^g_4=-3$									
4	2	$(0,-1)\times(0,1)\times(2,0)$										

TABLE II: The chiral and vector-like superfields, and their quantum numbers under the gauge symmetry $SU(4)_C \times SU(2)_L \times SU(2)_R \times USp(2)_1 \times USp(2)_2 \times USp(2)_3 \times USp(2)_4$.

	Quantum Number	Q_4	Q_{2L}	Q_{2R}	Field
ab	$3\times(4,\overline{2},1,1,1,1,1)$	1	-1	0	$F_L(Q_L,L_L)$
ac	$3\times(\overline{4},1,2,1,1,1,1)$	-1	0	1	$F_R(Q_R,L_R)$
$a1$	$1\times(4,1,1,2,1,1,1)$	1	0	0	
$a2$	$1\times(\overline{4},1,1,1,2,1,1)$	-1	0	0	
$b2$	$1\times(1,2,1,1,2,1,1)$	0	1	0	
$b4$	$3\times(1,\overline{2},1,1,1,1,2)$	0	-1	0	
$c1$	$1\times(1,1,\overline{2},2,1,1,1)$	0	0	-1	
$c3$	$3\times(1,1,2,1,1,2,1)$	0	0	1	
b_S	$2\times(1,3,1,1,1,1,1)$	0	2	0	T^i_L
b_A	$2\times(1,\overline{1},1,1,1,1,1)$	0	-2	0	S^i_L
c_S	$2\times(1,1,\overline{3},1,1,1,1)$	0	0	-2	T^i_R
c_A	$2\times(1,1,1,1,1,1,1)$	0	0	2	S^i_R
ab'	$3\times(4,2,1,1,1,1,1)$	1	1	0	
	$3\times(\overline{4},\overline{2},1,1,1,1,1)$	-1	-1	0	
ac'	$3\times(4,1,2,1,1,1,1)$	1		1	Φ_i
	$3\times(\overline{4},1,\overline{2},1,1,1,1)$	-1	0	-1	$\overline{\Phi}_i$
bc	$6\times(1,2,\overline{2},1,1,1,1)$	0	1	-1	$H^i_u,\ H^i_d$
	$6\times(1,\overline{2},2,1,1,1,1)$	0	-1	1	

The anomalies from three global $U(1)$s of $U(4)_C$, $U(2)_L$ and $U(2)_R$ are cancelled by the Green-Schwarz mechanism, and the gauge fields of these $U(1)$s obtain masses via the linear $B\wedge F$ couplings. Thus, the effective gauge symmetry is $SU(4)_C\times SU(2)_L\times SU(2)_R$. In order to break the gauge symmetry, on the first torus, we split the a stack of D6-branes into a_1 and a_2 stacks with 6 and 2 D6-branes, respectively, and split the c stack of D6-branes into c_1 and c_2 stacks with two D6-branes for each one. In this way, the gauge symmetry is further broken to $SU(3)_C\times SU(2)_L\times U(1)_{I_{3R}}\times U(1)_{B-L}$. Moreover, the $U(1)_{I_{3R}}\times U(1)_{B-L}$ gauge symmetry may be broken to $U(1)_Y$ by giving vacuum expectation values (VEVs) to the vector-like particles with the quantum numbers $(\mathbf{1},\mathbf{1},\mathbf{1/2},\mathbf{-1})$ and $(\mathbf{1},\mathbf{1},\mathbf{-1/2},\mathbf{1})$ under the $SU(3)_C\times SU(2)_L\times U(1)_{I_{3R}}\times U(1)_{B-L}$ gauge symmetry from a_2c_1' intersections [7, 10].

Since the gauge couplings in the Minimal Supersymmetric Standard Model (MSSM) are unified at the GUT scale $\sim 2.4\times10^{16}$ GeV, the additional exotic particles present in the model must necessarily become superheavy. To accomplish this it is first assumed that the $USp(2)_1$ and $USp(2)_2$ stacks of D6-branes lie on the top of each other on the first torus, so we have two pairs of vector-like particles with $USp(2)_1\times USp(2)_2$ quantum numbers $(2,2)$. These particles can break $USp(2)_1\times USp(2)_2$ down to the diagonal $USp(2)_{D12}$ near the string scale, and then states arising from intersections $a1$ and $a2$ may obtain vector-like masses close to the string scale. Moreover, we assume that the T^i_R and S^i_R obtain VEVs near the string scale, and their VEVs satisfy the D-flatness of $U(1)_R$. To preserve the D-flatness of $U(1)_L$, we assume that the VEVs of S^i_L is TeV scale. We also assume that there exist various suitable high-dimensional operators in the effective theory. With T^i_R and S^i_R, we can give the GUT-scale masses to the particles from the intersections $c1$, $c3$, and c_S via high-dimensional operators. The remaining states and adjoint chiral superfields may also obtain GUT-scale masses via high-dimensional operators by the Higgs mechanism and from strong dynamics since all of the $USp(2)_i$ have negative beta functions as shown in Table I [11]. To have one pair of light Higgs doublets, it is necessary to fine-tune the mixing parameters of the Higgs doublets. In particular, the μ term and the right-handed neutrino masses may be generated via the following high-dimensional operators

$$W \supset \frac{y^{ijkl}_{\mu}}{M_{\rm St}} S^i_L S^j_R H^k_u H^l_d + \frac{y^{mnkl}_{Nij}}{M^3_{\rm St}} T^m_R T^n_R \Phi_i \Phi_j F^k_R F^l_R \ , \tag{3}$$

where y^{ijkl}_{μ} and y^{mnkl}_{Nij} are Yukawa couplings, and $M_{\rm St}$ is the string scale. Thus, the μ term is TeV scale and the right-handed neutrino masses can be in the range 10^{10-14} GeV for $y^{ijkl}_{\mu}\sim1$ and $y^{mnkl}_{Nij}\sim10^{(-7)-(-3)}$.

Phenomenological Consequences $-$ In the string theory basis, we have the dilaton S, three Kähler moduli T^i, and three complex structure moduli U^i [12]. The U^i for the present model are

$$U^1 = 3i\ , \quad U^2 = i\ , \quad U^3 = -1+i\ . \tag{4}$$

The corresponding moduli s, t^i and u^i in the supergravity theory basis are related to the S, T^i and U^i moduli by [12]

$$\mathrm{Re}\,(s) \;=\; \frac{e^{-\phi_4}}{2\pi}\left(\frac{\sqrt{U_2^1\,U_2^2\,U_2^3}}{|U^1U^2U^3|}\right)\;,\;\mathrm{Re}(t^j) \;=\; \frac{i\alpha'}{T^j}\;,$$
$$\mathrm{Re}\,(u^j) \;=\; \frac{e^{-\phi_4}}{2\pi}\left(\sqrt{\frac{U_2^j}{U_2^k\,U_2^l}}\right)\left|\frac{U^k\,U^l}{U^j}\right|\;, \qquad (5)$$

where ϕ_4 is the four-dimensional dilaton, U_2^i is the imaginary part of U^i, and $j \neq k \neq l \neq j$.

The holomorphic gauge kinetic function for a generic P stack of D6-branes which does not lie on one of O6-planes, is given by [12]

$$f_P \;=\; \frac{1}{8}\left(2n_P^1\,n_P^2\,n_P^3\,s - n_P^1\,l_P^2\,l_P^3\,u^1\right.$$
$$\left.-n_P^2\,l_P^1\,l_P^3\,u^2 - 2n_P^3\,l_P^1\,l_P^2\,u^3\right)\;. \qquad (6)$$

Thus, the gauge couplings for $SU(4)_C$, $SU(2)_L$ and $SU(2)_R$ in our model are unified at the string scale. For simplicity, we negelect the little hierarchy between the string scale and the GUT scale, which may be explained via threshold corrections. Assuming the value of the unified gauge coupling in the MSSM, we obtain

$$e^{-\phi_4} = 20.1\;. \qquad (7)$$

Thus, the string scale is $\sim 2.1 \times 10^{17}$ GeV for $M_{\mathrm{St}} = \pi^{1/2}e^{\phi_4}M_{\mathrm{Pl}}$ where M_{Pl} is the reduced Planck scale.

The Kähler metric for the chiral superfields from the intersections of the P and Q stacks of D6-branes is [12]

$$\tilde{K} \supset e^{\phi_4+\gamma_E\sum_{i=1}^3\theta_{PQ}^i}\prod_{j=1}^3\left[\sqrt{\frac{\Gamma(1-\theta_{PQ}^i)}{\Gamma(\theta_{PQ}^i)}}(t^j+\bar{t}^j)^{-\theta_{PQ}^i}\right],$$

where γ_E is the Euler-Mascheroni constant, and θ_{PQ}^i is the suitable positive angle between the P and Q stacks of D6-branes on the i^{th} two torus in units of π [11], and can be written as a function of s, u^i, and the wrapping numbers for the P and Q stacks of D6-branes.

The Kähler metric for the vector-like chiral superfields from the intersections of the P and Q stacks of D6-branes that are parallel on the j^{th} two torus and intersect on the k^{th} and l^{th} two tori is given by [12]

$$\tilde{K} \supset \left[(s+\overline{s})(u^j+\overline{u}^j)(t^k+\overline{t}^k)(t^l+\overline{t}^l)\right]^{-1/2}\;. \qquad (8)$$

TABLE III: Supersymmetry breaking soft terms (in GeV) at the string scale.

M_1	M_2	M_3	m_{FL}	m_{FR}	m_H	A_Y
477.4	279.1	987.8	1047	524.7	451.7	732.6

For simplicity, we assume that only the F terms of the complex structure moduli u^i break supersymmetry and are parametrized as follows

$$F^{u^i} = \sqrt{3}m_{3/2}(u^i+\bar{u}^i)\Theta_i\;,\quad \text{for } i=1,2,3\;, \qquad (9)$$

where $m_{3/2}$ is the gravitino mass, and Θ_i are real numbers and satisfy $\sum_{i=1}^3|\Theta_i|^2 = 1$. Then, we can calculate the gaugino masses (M_i), the universal scalar masses m_{FL} and m_{FR} respectively for the left-handed and right-handed SM fermions, the universal scalar mass m_H for Higgs fields H_u^i and H_d^i, and the universal trilinear soft term A_Y at the string scale [13]. Choosing $m_{3/2} = 1100$ GeV, $\Theta_1 = -0.6$, $\Theta_2 = 0.293$, $\Theta_3 = 0.744$, $\mathrm{Re}t_1 = 1/6.6$, and $\mathrm{Re}t_2 = \mathrm{Re}t_3 = 0.5$, we obtain the string-scale supersymmetry breaking soft terms given in Table III. Using the code `SuSpect` [14], we calculate the low energy supersymmetric particle spectrum. An example for $\tan\beta = 46$ and positive μ is shown in Table IV. This spectrum is consistent with all the known experiments and can be tested at the LHC. Finally, using the code `MicrOMEGAs` [15], we obtain a dark matter density $\Omega h^2 = 0.117$ which is very close to the observed value.

TABLE IV: Low energy supersymmetric particles and their masses (in GeV).

h^0	H^0	A^0	$H^\pm$	$\tilde{g}$	$\chi_1^\pm$	$\chi_2^\pm$	χ_1^0	χ_2^0
121.3	1016	1017	1020	2192	219.3	1406	199.3	219.4
χ_3^0	χ_4^0	$\tilde{t}_1$	$\tilde{t}_2$	$\tilde{u}_1/\tilde{c}_1$	$\tilde{u}_2/\tilde{c}_2$	$\tilde{b}_1$	$\tilde{b}_2$	
1404	1405	1542	1912	1948	2144	1763	1915	
$\tilde{d}_1/\tilde{s}_1$	$\tilde{d}_2/\tilde{s}_2$	$\tilde{\tau}_1$	$\tilde{\tau}_2$	$\tilde{\nu}_\tau$	$\tilde{e}_1/\tilde{\mu}_1$	$\tilde{e}_2/\tilde{\mu}_2$	$\tilde{\nu}_e/\tilde{\nu}_\mu$	
1947	2146	234.4	1010	1000	550.2	1059	1056	

The SM Fermion Masses and Mixings – Because all the SM fermions and Higgs fields arise from the intersections on the first torus, we will only consider it for simplicity. The up-type quark mass matrix M^U at the GUT scale is [16]

$$c_0^U\begin{pmatrix} A^Uv_u^1+E^Uv_u^4 & B^Uv_u^3+F^Uv_u^6 & D^Uv_u^2+C^Uv_u^5 \\ C^Uv_u^3+D^Uv_u^6 & A^Uv_u^5+E^Uv_u^2 & B^Uv_u^1+F^Uv_u^4 \\ F^Uv_u^2+B^Uv_u^5 & C^Uv_u^1+D^Uv_u^4 & A^Uv_u^3+E^Uv_u^6 \end{pmatrix},$$

where $v_u^i = \langle H_u^i\rangle$, and c_0^U is a constant which includes the quantum corrections and the contributions to the Yukawa couplings from the second and third two tori. The theta functions A^U, B^U, C^U, D^U, E^U, and F^U are

$$A^U \equiv \vartheta\begin{bmatrix}\epsilon^{U1}\\ \phi^{(1)}\end{bmatrix}(\kappa^{(1)}),\;\; B^U \equiv \vartheta\begin{bmatrix}\epsilon^{U1}+\frac{1}{3}\\ \phi^{(1)}\end{bmatrix}(\kappa^{(1)})\;,$$
$$C^U \equiv \vartheta\begin{bmatrix}\epsilon^{U1}-\frac{1}{3}\\ \phi^{(1)}\end{bmatrix}(\kappa^{(1)}),\;\; D^U \equiv \vartheta\begin{bmatrix}\epsilon^{U1}+\frac{1}{6}\\ \phi^{(1)}\end{bmatrix}(\kappa^{(1)}),$$
$$E^U \equiv \vartheta\begin{bmatrix}\epsilon^{U1}+\frac{1}{2}\\ \phi^{(1)}\end{bmatrix}(\kappa^{(1)}),\;\; F^U \equiv \vartheta\begin{bmatrix}\epsilon^{U1}-\frac{1}{6}\\ \phi^{(1)}\end{bmatrix}(\kappa^{(1)}),$$

where

$$\epsilon^{U1} \equiv \frac{\epsilon_c^{U1} - \epsilon_b^{U1} - 2\epsilon_a^{U1}}{6}, \ \kappa^{(1)} \equiv \frac{6J^{(1)}}{\alpha'},$$
$$\phi^{(1)} = \quad \theta_c^{(1)} - \theta_b^{(1)} - 2\theta_a^{(1)}, \tag{10}$$

where ϵ_a^{U1}, ϵ_b^{U1} and ϵ_c^{U1} respectively are the shifts of a, b, and c stacks of D6-branes, $J^{(1)}$ is the Kähler modulus, and $\theta_a^{(1)}$, $\theta_b^{(1)}$ and $\theta_c^{(1)}$ are the Wilson line phases for the a, b, and c stacks on the first two torus, respectively.

At the GUT scale, the down-type quark mass matrix M^D is obtained from the above up-type quark mass matrix M^U by changing the upper index U and lower index u to D and d, respectively. The lepton mass matrix M^L is obtained from M^D by changing the upper index D to L.

To generate the suitable SM fermion masses and mixings at the GUT scale, we choose $\epsilon^{U1} = \epsilon^{L1} = 0$, $\epsilon^{D1} = 0.061$, and $\kappa^{(1)} = 39.6i$. And for Higgs VEVs, we choose $v_1^u = 0.000266$, $v_2^u = 0.236$, $v_3^u = 0.999$, $v_4^u = 0.981$, $v_5^u = 0.00481$, $v_6^u = 0.0345$, $v_1^d = 0.00224$, $v_2^d = 0$, $v_3^d = 1.58$, $v_4^d = 0$, $v_5^d = 0.0445$, and $v_6^d = 0.0001$. Then, with suitable c_0^U, c_0^D, and c_0^L, we obtain the SM fermion mass matrices at the GUT scale

$$M^U \simeq m_t \begin{pmatrix} 0.000266 & 0.00109 & 0.00747 \\ 0.00109 & 0.00481 & 0.0310 \\ 0.00747 & 0.0310 & 0.999 \end{pmatrix},$$
$$M^D \simeq m_b \begin{pmatrix} 0.00141 & 0.000025 & 4\times 10^{-6} \\ 0.000155 & 0.028 & 0.0 \\ 0.0 & 2.2\times 10^{-7} & 1 \end{pmatrix},$$
$$M^L \simeq m_\tau \begin{pmatrix} 0.00142 & 3.0\times 10^{-6} & 2.8\times 10^{-8} \\ 3.0\times 10^{-6} & 0.0282 & 1.4\times 10^{-9} \\ 2.8\times 10^{-8} & 1.4\times 10^{-9} & 1 \end{pmatrix}.$$

The above mass matrices can produce the correct quark masses and CKM mixings, and the correct τ lepton mass at the electroweak scale [17]. The electron mass is about 6.5 times larger that the expected value, while the muon mass is about 40% smaller. Similar to the GUTs [18], we have roughly the wrong fermion mass relation $m_e/m_\mu \simeq m_d/m_s$, and the correct electron and muon masses can be generated via high-dimensional operators [11]. Moreover, the suitable neutrino masses and mixings can be generated via the seesaw mechanism by choosing suitable Majorana mass matrix for the right-handed neutrinos.

Conclusions – We have briefly presented a three-family intersecting D6-brane model where the gauge symmetry can be broken down to the SM gauge symmetry and the gauge coupling unification can be realized at the string scale. We have calculated the supersymmetry breaking soft terms, and obtained the low energy supersymmetric particle spectrum within the reach of the LHC. Our model may also generate the observed dark matter density. Finally, we can explain the SM quark masses and CKM mixings, and the tau lepton mass. The neutrino masses and mixings may be generated via the seesaw mechanism as well.

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